QCD Factorization For B Decays To Two Light Pseudoscalars Including Chirally Enhanced Corrections *

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Abstract

Since b quark mass is not asymptotically large, chirally enhanced corrections which arise from twist-3 wave functions may be important in B decays. We thus evaluate the hadronic matrix elements with the emitted meson described by leading twist and twist-3 distribution amplitudes $\Phi_p(x)$. After summing over the four "vertex correction" diagrams, we obtain the results with infrared finiteness which shows that chirally enhanced corrections arise from $\Phi_p(x)$ can be consistently included in QCD factorization. We also briefly discuss the contributions from "hard spectator" diagrams.

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It is well known that two-body, non-leptonic charmless B decays are crucial for extracting CKM matrix elements. However, due to our ignorance on how to calculate hadronic matrix elements, we conventionally resort to Bjorken's color transparency argument [1] which lead to "naive factorization assumption",

$$\langle M_1 M_2|Q|B\rangle = \langle M_2|J_1|0\rangle\langle M_1|J_2|B\rangle, \tag{1}$$

This assumption makes the hadronic matrix elements scale-independent. Noting that Wilson Coefficients are scheme- and scale-dependent, the theoretical calculations on the branching ratios would then be scheme- and scale-dependent which is unacceptable. To save factorization hypothesis, a phenomenological parameter N_{eff} is introduced [4] which is commonly called generalized factorization. However this approach is not satisfactory yet because in principle N_{eff} is process dependent.

Recently, Beneke, Buchalla, Neubert and Sachrajda [2,3] proposed a promising method: in the heavy quark limit, they show that the emitted light meson can be described with leading twist-2 distribution amplitude, the infrared divergences of the hard-scattering amplitudes are canceled after summing over the four "vertex correction" diagrams (Fig.(a)-(d)), which is a one-loop demonstration of Bjorken's color transparency argument [1]. In the heavy quark limit, they show that the hadronic matrix elements can be expressed as [2]

$$\langle M_1 M_2|Q|B\rangle = \langle M_2|J_1|0\rangle\langle M_1|J_2|B\rangle \cdot [1 + \Sigma r_n \alpha_s^n + \mathcal{O}(\Lambda_{QCD}/m_b)]. \tag{2}$$

If power corrections in $1/m_b$ can be safely neglected, then everything is perfect. At the zero order of α_s, it would come back to "naive factorization", and at the higher order of α_s, the corrections can be systematically calculated in Perturbative QCD, which means that the decay amplitudes of B meson can be computed from first principles, and the necessary input are heavy-to-light form factors and light-cone distribution amplitudes. But in the real world, bottom quark mass is not asymptotically large (but 4.8 GeV), and numerically power suppression may fail in some cases. An obvious and possibly the most important case is chirally enhanced power corrections. As pointed out in ref [2], numerically the enhanced factor $r_\chi = \frac{2m_\pi^2}{m_b(m_u+m_d)} \simeq 1.18$ which makes the power suppression completely fail. This parameter is multiplied by a_6 and a_8, where a_6 is very important numerically in penguin-dominated B decays. So an evaluation of the hadronic matrix elements including chirally enhanced corrections may be phenomenologically or numerically important. In this letter, we shall examine this problem in some detail.

Chirally enhanced corrections arise from twist-3 light-cone distribution amplitudes, generally called $\Phi_p(x)$ and $\Phi_\sigma(x)$. For light pseudoscalar mesons, they are defined as [6]

$$\langle P(p')|\bar{q}(y) i\gamma_5 q(x)|0\rangle = f_p \mu_p \int_0^1 du \;\; e^{i(up'\cdot y + \bar{u}p'\cdot x)} \phi_p(u), \tag{3}$$

$$\langle P(p')|\bar{q}(y)\sigma_{\mu\nu}\gamma_5 q(x)|0\rangle = i f_p \mu_p (p'_\mu z_\nu - p'_\nu z_\mu) \int_0^1 du \;\; e^{i(up'\cdot y + \bar{u}p'\cdot x)} \frac{\phi_\sigma(u)}{6}, \tag{4}$$

where $\mu_p = \frac{M_p^2}{m_u+m_d}$, $z = y - x$. We notice that in Ref. [2] color transparency is demonstrated in one-loop level in the heavy quark limit. If we want to include chirally enhanced corrections consistently, we should describe the emitted light meson with leading twist-2 and twist-3 distribution amplitudes, which means that we should show the infrared finiteness using

twist-3 distribution amplitudes after summing over the "vertex correction" diagrams. In this paper, we shall restrict ourselves to $\Phi_p(x)$ while postpone the discussion of $\Phi_\sigma(x)$ to ref [7] because of the complicated derivation in proving the infrared finiteness of the "vertex correction" diagrams using $\Phi_\sigma(x)$.

We notice that in Ref [8], the authors have used twist-3 distribution amplitude $\Phi_p(x)$ to calculate the strong penguin corrections (Fig.(e)-(f)). The difference of our work from ref [8] is that we calculate "vertex correction" diagrams and show the infrared finiteness of a_6 and a_8 at the order of α_s.

In the following, we take leading twist and twist-3 wave functions $\Phi_p(x)$ to describe the emitted light meson, and we will show the infrared finiteness under this approach.

The $|\Delta B| = 1$ effective Hamiltonian is given by [9]

$$\mathcal{H}_{eff} = \frac{G_F}{\sqrt{2}} \left[\sum_{q=u,c} v_q \left(C_1(\mu) Q_1^q(\mu) + C_2(\mu) Q_2^q(\mu) + \sum_{k=3}^{10} C_k(\mu) Q_k(\mu) \right) - v_t (C_{7\gamma} Q_{7\gamma} + C_{8G} Q_{8G}) \right] + h.c., \tag{5}$$

where $v_q = V_{qb} V_{qd}^*$(for $b \to d$ transition) or $v_q = V_{qb} V_{qs}^*$(for $b \to s$ transition) and $C_i(\mu)$ are Wilson coefficients which have been evaluated to next-to-leading order approximation. The four-quark operators Q_i are

$$\begin{array}{ll}
Q_1^u = (\bar{u}_\alpha b_\alpha)_{V-A} (\bar{q}_\beta u_\beta)_{V-A} & Q_1^c = (\bar{c}_\alpha b_\alpha)_{V-A} (\bar{q}_\beta c_\beta)_{V-A} \\
Q_2^u = (\bar{u}_\alpha u_\alpha)_{V-A} (\bar{q}_\beta b_\beta)_{V-A} & Q_2^c = (\bar{c}_\alpha c_\alpha)_{V-A} (\bar{q}_\beta b_\beta)_{V-A} \\
Q_3 = (\bar{q}_\alpha b_\alpha)_{V-A} \sum\limits_{q'} (\bar{q}'_\beta q'_\beta)_{V-A} & Q_4 = \sum\limits_{q'} (\bar{q}_\beta q'_\beta)_{V-A} (\bar{q}'_\alpha b_\alpha)_{V-A} \\
Q_5 = (\bar{q}_\alpha b_\alpha)_{V-A} \sum\limits_{q'} (\bar{q}'_\beta q'_\beta)_{V+A} & Q_6 = -2 \sum\limits_{q'} (\bar{q}_\beta q'_\beta)_{S+P} (\bar{q}'_\alpha b_\alpha)_{S-P} \\
Q_7 = \frac{3}{2} (\bar{q}_\alpha b_\alpha)_{V-A} \sum\limits_{q'} e_{q'} (\bar{q}'_\beta q'_\beta)_{V+A} & Q_8 = -3 \sum\limits_{q'} e_{q'} (\bar{q}_\beta q'_\beta)_{S+P} (\bar{q}'_\alpha b_\alpha)_{S-P} \\
Q_9 = \frac{3}{2} (\bar{q}_\alpha b_\alpha)_{V-A} \sum\limits_{q'} e_{q'} (\bar{q}'_\beta q'_\beta)_{V-A} & Q_{10} = \frac{3}{2} \sum\limits_{q'} e_{q'} (\bar{q}_\beta q'_\beta)_{V-A} (\bar{q}'_\alpha b_\alpha)_{V-A}
\end{array} \tag{6}$$

and

$$Q_{7\gamma} = \frac{e}{8\pi^2} m_b \bar{q}_\alpha \sigma^{\mu\nu} (1+\gamma_5) b_\alpha F_{\mu\nu}, \quad Q_{8G} = \frac{g}{8\pi^2} m_b \bar{q}_\alpha \sigma^{\mu\nu} t^a_{\alpha\beta} b_\beta G^a_{\mu\nu}, \quad (q = d \text{ or } s). \tag{7}$$

The amplitude of the decays of B to two light pseudoscalar mesons in QCD factorization can be written as:

$$A(B \to M_1 M_2) = \frac{G_F}{\sqrt{2}} \sum_{p=u,c} \sum_{i=1,10} v_p a_i^p \langle M_1 M_2 | Q_i | B \rangle_F, \tag{8}$$

where v_p is CKM factor and $\langle M_1 M_2 | Q_i | B \rangle_F$ is the factorized matrix elements.

We calculate QCD coefficients a_i^p with the emitted mesons described by light-cone distribution amplitudes. For instance, let we consider the "vertex correction" diagrams. Twist-3 distribution amplitude $\Phi_p(x)$ makes no contribution when considering $(V-A) \otimes (V-A)$ and $(S+P) \otimes (S-P)$ currents because of their Lorentz structure. As to $(V-A) \otimes (V+A)$ current, there is some subtlety in regularizing the infrared divergences. If we use dimension regularization, the infrared finiteness would not hold for the case of $\Phi_p(x)$ after summing

over those four "vertex correction" diagrams. That is because wave functions are defined in 4-dimensions, it may be unconsistent to extend its usage directly to D-dimensions. Thus we assign a virtual mass to the gluon propagator and regularize the infrared integrals in four dimensions. Then the "vertex correction" contributions of $(V-A)\otimes(V+A)$ current to $(S+P)\otimes(S-P)$ operator is infrared finite:

$$V=\frac{\alpha_s}{4\pi}\frac{C_F}{N}\int_0^1 dx\Phi_p(x)\{i\pi\log\frac{x}{\bar{x}}-\frac{1+x}{x}\log\bar{x}+\frac{1+\bar{x}}{\bar{x}}\log x-Li_2(-\frac{\bar{x}}{x})+Li_2(-\frac{x}{\bar{x}})\}, \quad (9)$$

where $\bar{x}=1-x$ and $Li_2(x)$ is dilogarithm function. This means that we can consistently include twist-3 distribution amplitude $\Phi_p(x)$ in the framework of QCD factorization. The explicit expressions of a_i^p for $i=1$ to 10 (for symmetric light-cone distribution amplitudes of light pseudoscalar mesons) are obtained as:

$$a_1^u=C_1+\frac{C_2}{N}+\frac{\alpha_s}{4\pi}\frac{C_F}{N}C_2F, \quad (10)$$

$$a_2^u=C_2+\frac{C_1}{N}+\frac{\alpha_s}{4\pi}\frac{C_F}{N}C_1F, \quad (11)$$

$$a_3=C_3+\frac{C_4}{N}+\frac{\alpha_s}{4\pi}\frac{C_F}{N}C_4F, \quad (12)$$

$$\begin{aligned}a_4^p=C_4+\frac{C_3}{N}-\frac{\alpha_s}{4\pi}\frac{C_F}{N}\Big\{&(\frac{4}{3}C_1+\frac{44}{3}C_3+\frac{4f}{3}(C_4+C_6))\ln\frac{\mu}{m_b}\\&+(G_{M_2}(s_p)-\frac{2}{3})C_1+(G_{M_2}(0)+G_{M_2}(1)-f^I_{M_2}-f^{II}_{M_2}+\frac{50}{3})C_3\\&-\frac{2f}{3}C_4+(3G_{M_2}(0)+G_{M_2}(s_c)+G_{M_2}(1))(C_4+C_6)+G_{M_2,8}C_8\Big\},\end{aligned} \quad (13)$$

$$a_5=C_5+\frac{C_6}{N}+\frac{\alpha_s}{4\pi}\frac{C_F}{N}C_6(-F-12), \quad (14)$$

$$\begin{aligned}a_6^p=C_6+\frac{C_5}{N}-\frac{\alpha_s}{4\pi}\frac{C_F}{N}6C_5-\frac{\alpha_s}{4\pi}\frac{C_F}{N}\Big\{&(C_1+2C_3+f(C_4+C_6))\ln\frac{\mu}{mb}\\&+(G'_{M_2}(s_p)-\frac{7}{12})C_1+(G'_{M_2}(0)+G'_{M_2}(1)-\frac{7}{6})C_3-\frac{7}{12}fC_4\\&+(3G'_{M_2}(0)+G'_{M_2}(s_c)+G'_{M_2}(1))(C_4+C_6)-\frac{f}{12}C_6+\frac{3}{2}C_{8G}\Big\},\end{aligned} \quad (15)$$

$$a_7=C_7+\frac{C_8}{N}+\frac{\alpha_s}{4\pi}\frac{C_F}{N}C_8(-F-12), \quad (16)$$

$$\begin{aligned}a_8=C_8+\frac{C_7}{N}&-\frac{\alpha_s}{4\pi}\frac{C_F}{N}6C_7\\&+\frac{\alpha_{em}}{9\pi}\Big\{-((\frac{C_1}{N}+C_2)+\frac{1}{2}(\frac{C_4}{N}+C_3)+\frac{1}{2}(\frac{C_6}{N}+C_5))\ln\frac{\mu}{mb}\\&+(\frac{7}{12}-G'_{M_2}(s_p))(\frac{C_1}{N}+C_2)+(\frac{7}{24}-G'_{M_2}(s_c)+\frac{1}{2}G'_{M_2}(1))(\frac{C_4}{N}+C_3)\\&+(\frac{1}{24}-G'_{M_2}(s_c)+\frac{1}{2}G'_{M_2}(1))(\frac{C_6}{N}+C_5)-\frac{3}{4}C_{7\gamma}\Big\},\end{aligned} \quad (17)$$

$$a_9=C_9+\frac{C_{10}}{N}+\frac{\alpha_s}{4\pi}\frac{C_F}{N}C_{10}F, \quad (18)$$

$$
\begin{aligned}
a_{10}^p = C_{10} + \frac{C_9}{N} + \frac{\alpha_s}{4\pi}\frac{C_F}{N}C_9 F \\
+\frac{\alpha_{em}}{9\pi}\Big\{(-\frac{2}{3}(2(C_2+\frac{C_1}{N})+(C_3+\frac{C_4}{N})+(C_5+\frac{C_6}{N}))\ln\frac{\mu}{m_b} \\
+(\frac{2}{3}-G_{M_2}(s_p))(C_2+\frac{C_1}{N})+(\frac{1}{3}-G_{M_2}(s_c)+\frac{G_{M_2}(1)}{2})(C_3+\frac{C_4}{N}) \\
+(-G_{M_2}(s_c)+\frac{G_{M_2}(1)}{2})(C_5+\frac{C_6}{N})-\frac{1}{2}C_{7\gamma}G_{M_2,8}\Big\}.
\end{aligned}
\tag{19}
$$

Here $N = 3$ ($f = 5$) is the number of color (flavor), $C_F = \frac{N^2-1}{2N}$ is the factor of color, $s_p = m_p^2/m_b^2$ for $p = u, c$ and we define the symbols in the above expressions as:(most of them are as the same as Beneke's except for $G'_{M_2}(s)$ and $G'(s,x)$)

$$F = -12\ln\frac{\mu}{m_b} - 18 + f^I_{M_2} + f^{II}_{M_2}, \tag{20}$$

$$f^I_{M_2} = \int_0^1 dx\ g(x)\Phi_{M_2}(x),\ \ G_{M_2,8} = \int_0^1 dx\ G_8(x)\Phi_{M_2}(x), \tag{21}$$

$$G_{M_2}(s) = \int_0^1 dx\ G(s,x)\Phi_{M_2}(x), \tag{22}$$

$$G'_{M_2}(s) = \int_0^1 dx\ G'(s,x)\Phi^p_{M_2}(x), \tag{23}$$

here $\Phi_{M_2}(x)(\Phi^p_{M_2}(x))$ is leading twist (twist-3) wave function of the emitted meson M_2, and the hard-scattering functions are

$$g(x) = 3\frac{1-2x}{1-x}\ln x - 3i\pi, \quad G_8(x) = \frac{2}{1-x}, \tag{24}$$

$$G(s,x) = -4\int_0^1 du\ u(1-u)\ln(s-u(1-u)(1-x)-i\epsilon), \tag{25}$$

$$G'(s,x) = \frac{3}{4}G(s,x). \tag{26}$$

As to $f^{II}_{M_2}$ which labels the contributions from the hard spectator scattering diagrams (Fig.(g)-(h)), We take the wave function of B meson as $\gamma_5(\not{P}_B - M_B)\Phi_B(\xi)$ and find that, when considering twist-3 $\Phi_p(x)$ distribution contributions, the hard spectator scattering contributions are proportional to :

(i) for the case of $(V-A)\otimes(V-A)$ and $(S+P)\otimes(S-P)$ currents ,

$$\int_0^1 d\xi\ \frac{\Phi_B(\xi)}{\xi}\int_0^1 dy\ \frac{\Phi_{M_1}(y)}{y}\int_0^1 dx\ \left\{\frac{\Phi_{M_2}(x)}{x} + \frac{\Phi^p_{M_2}(x)\frac{\mu_{M_2}}{M_B}(x-(1-x))}{x(1-x)}\right\}; \tag{27}$$

(ii) for the case of $(V-A)\otimes(V+A)$ currents,

$$\int_0^1 d\xi \ \frac{\Phi_B(\xi)}{\xi} \int_0^1 dy \ \frac{\mu_{M_1}}{M_B} \frac{\Phi^p_{M_1}(y)}{y} \int_0^1 dx \ \frac{\Phi^p_{M_2}(x) \frac{\mu_{M_2}}{M_B} (x-(1-x))}{x(1-x)}. \tag{28}$$

Under the symmetric distributions of final state light mesons, the logarithmically divergent integrals are canceled after summing over two hard spectator scattering diagrams. As a result there is no hard spectator scattering contributions to a_6 and a_8, the contribution of hard spectator scattering to other a_i^p is as the same as ref [2,10]:

$$f^{II}_{M_2} = \frac{4\pi^2}{N} \frac{f_{M_1} f_B}{F_+^{B\to M_1}(0) m_B^2} \int_0^1 d\xi \ \frac{\Phi_B(\xi)}{\xi} \int_0^1 dx \ \frac{\Phi_{M_1}(x)}{x} \int_0^1 dy \ \frac{\Phi_{M_2}(y)}{y}. \tag{29}$$

In ref [5], the authors have discussed the contributions of asymmetric distributions and find that numerically this effect is very small. In our case, when considering the asymmetric distributions, there would appear divergent integrals, but in this case the asymmetric distribution corrections would also be small if we parametrized the divergent integrals as an unknown parameter(as what have done in ref [5]) and could be safely neglected.

We notice that the above approach of evaluating "hard spectator" contribution is naive. For instance, the scale of "hard spectator" contribution should be different from the "vertex correction" contribution. While it seems reasonable, for the "vertex correction" diagrams, to take the scale $\mu \sim \mathcal{O}(m_b)$ to avoid large logarithm $\alpha_s \log \frac{\mu}{m_b}$, a natural choice of the scale of "hard spectator" contribution may be around $\mathcal{O}(1\ GeV)$ because the average momentum square of the exchanged gluon is about $1\ GeV^2$. Another disturbing feature of "hard spectator" contribution is that, as have been pointed out in ref [5], when including the contribution of Φ_σ, there would appear divergent integral $\int_0^1 dx \frac{1}{x}$ even if the symmetric distribution amplitude is applied. This divergent integral implies that the dominant contribution comes from the end-point region, or in another word, it is dominated by soft gluon exchange. However the transverse momentum may not be omitted in the end-point region, if so, the corresponding divergent integral would then changed to:

$$\int dx d^2 k_T \frac{\phi(x, k_T)}{x\xi m_b^2 + k_T^2}. \tag{30}$$

As an illustration, we do not consider the k_T dependence of wave functions (though it is certainly not a good approximation), then the above integral proportions to:

$$\int \frac{dx dk_T^2}{x\xi m_b^2 + k_T^2} \propto \int \frac{dx dy}{x+y}. \tag{31}$$

The above integration is convergence now, furthermore it is not dominated by end-point contribution. This illustrates that the treatment of "hard spectator" diagrams may need further discussions.

In summary, we consider some chirally enhanced corrections arise from twist-3 distribution amplitudes $\Phi_p(x)$. To include chirally enhanced corrections consistently in QCD factorization, we describe the emitted light meson with leading twist and twist-3 $\Phi_p(x)$ distribution amplitudes and show the infrared finiteness of "vertex correction" diagrams (fig.(a)-(d)). We also briefly discuss the disturbing "hard spectator" contributions.

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TABLES

QCD Coefficients	$\mu = 5.0\ GeV$ NLO	LO	$\mu = 2.5\ GeV$ NLO	LO
a_1^u	$1.043 + 0.012i$	1.019	$1.067 + 0.025i$	1.039
a_2^u	$0.034 - 0.075i$	0.184	$-0.011 - 0.104i$	0.104
a_3	$0.007 + 0.002i$	0.003	$0.010 + 0.004i$	0.005
a_4^u	$-0.030 - 0.014i$	-0.029	$-0.033 - 0.018i$	-0.041
a_4^c	$-0.035 - 0.006i$	-0.029	$-0.040 - 0.006i$	-0.041
a_5	$-0.007 - 0.003i$	-0.005	$-0.009 - 0.006i$	-0.010
$r_\chi a_6^u$	$-0.047 - 0.003i$	-0.044	$-0.050 - 0.003i$	-0.052
$r_\chi a_6^c$	$-0.049 - 0.006i$	-0.044	$-0.053 - 0.006i$	-0.052
$a_7 \times 10^5$	$11.1 + 2.7i$	9.1	$2.8 + 5.5i$	3.6
$r_\chi a_8^u \times 10^5$	$49.0 - 3.0i$	45.3	$57.0 - 1.2i$	49.2
$r_\chi a_8^c \times 10^5$	$47.3 - 4.6i$	45.3	$56.3 - 1.9i$	49.2
$a_9 \times 10^5$	$-939.8 - 13.3i$	-913.3	$-969.6 - 25.4i$	-941.4
$a_{10}^u \times 10^5$	$22.4 + 57.8i$	-135.7	$68.1 + 88.9i$	-68.0
$a_{10}^c \times 10^5$	$19.1 + 63.0i$	-135.7	$66.2 + 91.9i$	-68.0

TABLE I. The QCD coefficients $a_i^p(\pi\pi)$ at NLO and LO for the renormalization scales at $\mu = 5\ GeV$ and $\mu = 2.5\ GeV$

FIGURES

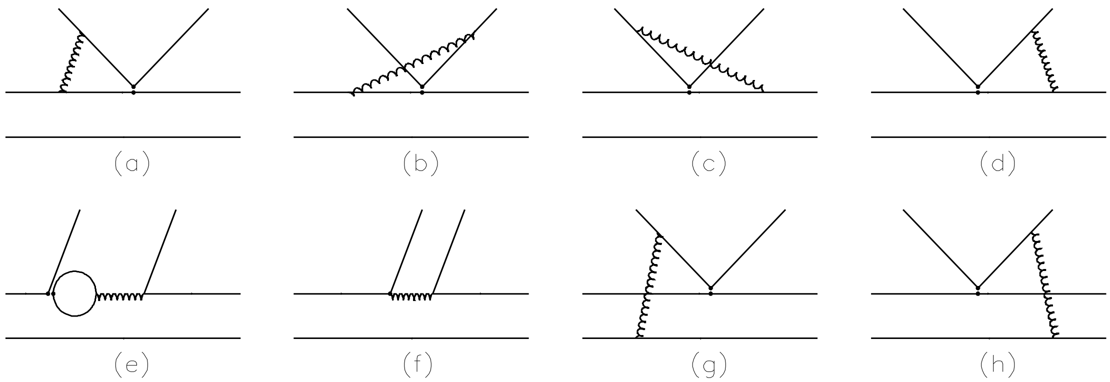

FIG. 1. Order of α_s corrections to hard-scattering kernels. The upward quark lines represent the ejected quark pairs from b quark weak decays.